

# MAGNETIC CATALYSIS AND INVERSE MAGNETIC CATALYSIS IN NONLOCAL CHIRAL QUARK MODELS

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## Abstract

We study the behavior of strongly interacting matter under an external constant magnetic field in the context of nonlocal chiral quark models within the mean field approximation. We find that at zero temperature the behavior of the quark condensates shows the expected magnetic catalysis effect, our predictions being in good quantitative agreement with lattice QCD results. On the other hand, in contrast to what happens in the standard local Nambu–Jona-Lasinio model, when the analysis is extended to the case of finite temperature our results show that nonlocal models naturally lead to the Inverse Magnetic Catalysis effect.

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Over the last few years, the understanding of the behavior of strongly interacting matter under extremely intense magnetic fields has attracted increasing attention, due to its relevance for various subjects such as the physics of compact objects like magnetars [1], the analysis of heavy ion collisions at very high energies [2] and the study of the first phases of the Universe [3]. Consequently, considerable work has been devoted to studying the structure of the QCD phase diagram in the presence of an external magnetic field (see Refs. [4–6] for recent reviews). On the basis of the results arising from most low-energy effective models of QCD it was generally expected that, at zero chemical potential, the magnetic field would lead to an enhancement of the chiral condensate (“magnetic catalysis”), independently of the temperature of the system. However, lattice QCD (LQCD) calculations carried out with physical pion masses [7, 8] show that, whereas at low temperatures one finds indeed such an enhancement, the situation is quite different close to the critical chiral restoration temperature: in that region light quark condensates exhibit a nonmonotonic behavior as functions of the external magnetic field, which results in a decrease of the transition temperature when the magnetic field is increased. This effect is known as inverse magnetic catalysis (IMC). Although many scenarios have been considered in the last few years to account for the IMC [9–29], the mechanism behind this effect is not yet fully understood. With this motivation, in this work we study the behavior of strongly interacting matter under an external magnetic field in the framework of nonlocal chiral quark models. These theories are proposed as a sort of nonlocal extensions of the well-known Nambu–Jona-Lasinio (NJL) model, intending to go a step further toward a more realistic effective approach to QCD. In fact, nonlocality arises naturally in the context of successful descriptions of low-energy quark dynamics [30, 31], and it has been shown [32] that nonlocal models can lead to a momentum dependence in quark propagators that is consistent with LQCD results. Another advantage of these models is that the effective interaction is finite to all orders in the loop expansion, and therefore there is no need to introduce extra cutoffs [33]. Moreover, in this framework it is possible to obtain an adequate description of the properties of light mesons at both zero and finite temperature/density [32, 34–43]. A previous attempt of considering the effect of an external magnetic field within these models was done in Ref. [44]. In that work the magnetic field was introduced by using a simplified extension of the method usually followed in the local NJL model, and no signal of IMC was found. In the present article we concentrate on the analysis of nonlocal quark models with separable interactions, including

a coupling to a uniform magnetic field. We address the problem by following a more rigorous procedure based on the Ritus eigenfunction method [45], which allows us to properly obtain the corresponding mean field action and to derive the associated gap equation. Then we solve this equation numerically for different values of the external magnetic field, considering the case of systems at both zero and finite temperature. We find that at zero temperature the behavior of the quark condensates shows the expected magnetic catalysis effect, our predictions being in good quantitative agreement with LQCD results. On the other hand, in contrast to what happens in the standard local Nambu–Jona-Lasinio model, when the analysis is extended to the case of finite temperature our results show that nonlocal models naturally lead to the IMC effect already at the mean field level.

### *Theoretical formalism*

We begin by stating the Euclidean action for a simple nonlocal chiral quark model that includes two light flavors,

$$S_E = \int d^4x \left\{ \bar{\psi}(x) (-i\not{D} + m_c) \psi(x) - \frac{G}{2} j_a(x) j_a(x) \right\} . \quad (1)$$

Here  $m_c$  is the current quark mass, which is assumed to be equal for  $u$  and  $d$  quarks. The nonlocal currents  $j_a(x)$  are given by

$$j_a(x) = \int d^4z \mathcal{G}(z) \bar{\psi}(x + \frac{z}{2}) \Gamma_a \psi(x - \frac{z}{2}) , \quad (2)$$

where  $\Gamma_a = (\mathbb{1}, i\gamma_5 \vec{\tau})$ , and the function  $\mathcal{G}(z)$  is a nonlocal form factor that characterizes the effective interaction. Since we are interested in studying the influence of a magnetic field, we introduce in the effective action Eq. (1) a coupling to an external electromagnetic gauge field  $\mathcal{A}_\mu$ . For a local theory this can be done by performing the replacement  $\partial_\mu \rightarrow \partial_\mu - i \hat{Q} \mathcal{A}_\mu(x)$ , where  $\hat{Q} = \text{diag}(q_u, q_d)$ , with  $q_u = 2e/3$ ,  $q_d = -e/3$ , is the electromagnetic quark charge operator. In the case of the nonlocal model under consideration the situation is more complicated since the inclusion of gauge interactions implies a change not only in the kinetic terms of the Lagrangian but also in the nonlocal currents in Eq. (2). One has

$$\psi(x - z/2) \rightarrow W(x, x - z/2) \psi(x - z/2) , \quad (3)$$

and a related change holds for  $\bar{\psi}(x + z/2)$  [32, 40, 43]. Here the function  $W(s, t)$  is defined by

$$W(s, t) = \text{P exp} \left[ -i \hat{Q} \int_s^t dr_\mu \mathcal{A}_\mu(r) \right] , \quad (4)$$

where  $r$  runs over an arbitrary path connecting  $s$  with  $t$ . As is usually done, we take it to be a straight line path.

To proceed we bosonize the fermionic theory, introducing scalar and pseudoscalar fields  $\sigma(x)$  and  $\vec{\pi}(x)$  and integrating out the fermion fields. The bosonized action can be written as [32, 43]

$$S_{\text{bos}} = -\ln \det \mathcal{D} + \frac{1}{2G} \int d^4x \left[ \sigma(x)\sigma(x) + \vec{\pi}(x) \cdot \vec{\pi}(x) \right], \quad (5)$$

where

$$\begin{aligned} \mathcal{D} \left( x + \frac{z}{2}, x - \frac{z}{2} \right) = & \gamma_0 W \left( x + \frac{z}{2}, x \right) \gamma_0 \left[ \delta^{(4)}(z) (-i\not{D} + m_c) + \right. \\ & \left. \mathcal{G}(z) [\sigma(x) + i\vec{\tau} \cdot \vec{\pi}(x)] \right] W \left( x, x - \frac{z}{2} \right). \end{aligned} \quad (6)$$

Let us consider the particular case of a constant and homogenous magnetic field orientated along the 3-axis. Choosing the Landau gauge, the corresponding gauge field is given by  $\mathcal{A}_\mu = B x_1 \delta_{\mu 2}$ . Next, we assume that the field  $\sigma$  has a nontrivial translational invariant mean field value  $\bar{\sigma}$ , while the mean field values of pseudoscalar fields  $\pi_i$  are zero. It should be stressed at this point that the assumption that  $\bar{\sigma}$  is independent of  $x$  does not imply that the resulting quark propagator will be translational invariant. In fact, as discussed below, one can show that such an invariance is broken by the appearance of the usual Schwinger phase. Our assumption just states that the deviations from translational invariance that are inherent to the magnetic field are not affected by the dynamics of the theory. In this way, within the mean field approximation (MFA) we get

$$\begin{aligned} \mathcal{D}^{\text{MFA}}(x, x') = & \delta^{(4)}(x - x') \left( -i\not{D} - \hat{Q} B x_1 \gamma_2 + m_c \right) + \\ & \bar{\sigma} \mathcal{G}(x - x') \exp \left[ \frac{i}{2} \hat{Q} B (x_2 - x'_2) (x_1 + x'_1) \right]. \end{aligned} \quad (7)$$

At this stage it is convenient to follow the Ritus eigenfunction method [45]. Thus, we introduce the function

$$\mathcal{D}_{p,p'}^{\text{MFA}} = \int d^4x d^4x' \bar{\mathbb{E}}_p(x) \mathcal{D}^{\text{MFA}}(x, x') \mathbb{E}_{p'}(x'), \quad (8)$$

where  $\mathbb{E}_p$  are the usual Ritus matrices, with  $p = (k, p_2, p_3, p_4)$ . The r.h.s. of Eq. (8) can be worked out, and after some calculation one arrives at a relatively compact expression for

$\mathcal{D}_{p,p'}^{\text{MFA}}$ , which is shown to be diagonal in flavor space. For each flavor  $f = u, d$  one has

$$\mathcal{D}_{p,p'}^{\text{MFA},f} = (2\pi)^4 \delta_{kk'} \delta(p_2 - p_2') \delta(p_3 - p_3') \delta(p_4 - p_4') \times \left[ [\mathbb{I} + \delta_{k0}(\Delta^{s_f} - \mathbb{I})] (-s_f \sqrt{2k|q_f B|} \gamma_2 + p_3 \gamma_3 + p_4 \gamma_4) + \sum_{\lambda=-1,1} \Delta^\lambda M_{\bar{p},k}^{\lambda,f} \right], \quad (9)$$

where we have defined  $s_f = \text{sign}(q_f B)$  and  $\Delta^\lambda = \text{diag}(\delta_{1\lambda}, \delta_{-1\lambda}, \delta_{1\lambda}, \delta_{-1\lambda})$ , whereas  $M_{\bar{p},k}^{\lambda,f}$  is given by

$$M_{\bar{p},k}^{\lambda,f} = (-1)^{k-\frac{1-\lambda s_f}{2}} \int_0^\infty dr r \exp(-r^2/2) \left[ m_c + \bar{\sigma} g \left( \frac{|q_f B|}{2} r^2 + \bar{p}^2 \right) \right] L_{k-\frac{1-\lambda s_f}{2}}(r^2). \quad (10)$$

Here  $g(p^2)$  stands for the Fourier transform of  $\mathcal{G}(z)$ ,  $\bar{p} = (p_3, p_4)$  is a two-dimensional vector and  $L_n(x)$  are the Laguerre polynomials. We use the standard convention  $L_{-1}(x) = 0$ , hence  $M_{\bar{p},0}^{-s_f,f} = 0$ . From Eq. (9) we finally find that the MFA action per unit volume can be expressed as

$$\frac{S_{\text{bos}}^{\text{MFA}}}{V^{(4)}} = \frac{\bar{\sigma}^2}{2G} - N_c \sum_{f=u,d} \frac{|q_f B|}{2\pi} \int \frac{d^2 \bar{p}}{(2\pi)^2} \left\{ \ln \left[ \bar{p}^2 + \left( M_{\bar{p},0}^{s_f,f} \right)^2 \right] + \sum_{k=1}^\infty \ln \left[ \left( 2k|q_f B| + \bar{p}^2 + M_{\bar{p},k}^{-1,f} M_{\bar{p},k}^{+1,f} \right)^2 + \bar{p}^2 \left( M_{\bar{p},k}^{+1,f} - M_{\bar{p},k}^{-1,f} \right)^2 \right] \right\}. \quad (11)$$

The corresponding gap equation can be now easily found by minimizing this expression with respect to  $\bar{\sigma}$ . It is worth mentioning that this gap equation can be also obtained using the Schwinger-Dyson formalism for the quark propagator discussed in e.g. Refs. [46–48]. Actually, it turns out that the two point function in Eq. (9) can be casted into a form similar to that given in Ref. [47]. Thus, using the analysis discussed in that work, one can show that the associated quark propagator in coordinate space can be written as the product of the exponential of a Schwinger phase (which breaks translational invariance) times a translational invariant function.

The above results can be now extended to finite temperature using the Matsubara formalism. This amounts to performing the replacement

$$\int \frac{d^2 \bar{p}}{(2\pi)^2} F(\bar{p}^2) \rightarrow T \sum_{n=-\infty}^\infty \int \frac{dp_3}{2\pi} F(\bar{p}_n^2), \quad (12)$$

where  $\bar{p}_n = (p_3, \omega_n)$ ,  $\omega_n = (2n+1)\pi T$  being the Matsubara frequencies for fermionic modes. In this way one can easily obtain the MFA finite temperature thermodynamical potential  $\Omega^{\text{MFA}}$ , as well as the related gap equation. Given  $\Omega^{\text{MFA}}$ , the magnetic field dependent quark

condensate for each flavor can be calculated by taking the derivative with respect to the corresponding current quark mass. This leads to

$$\langle \bar{q}_f q_f \rangle_{B,T} = -\frac{N_c |q_f B| T}{\pi} \int \frac{dp_3}{2\pi} \sum_{k=0}^{\infty} \sum_{n=-\infty}^{\infty} \frac{M_{\bar{p}_n,k}^{-1,f} \left[ \bar{p}_n^2 + 2k|q_f B| + M_{\bar{p}_n,k}^{+1,f^2} \right] + (+ \leftrightarrow -)}{\left( 2k|q_f B| + \bar{p}_n^2 + M_{\bar{p}_n,k}^{-1,f} M_{\bar{p}_n,k}^{+1,f} \right)^2 + \bar{p}_n^2 \left( M_{\bar{p}_n,k}^{+1,f} - M_{\bar{p}_n,k}^{-1,f} \right)^2} . \quad (13)$$

As it is usually found in the context of nonlocal models [39], this expression turns out to be divergent beyond the chiral limit. We obtain a regularized condensate by subtracting the corresponding expression in the absence of quark-quark interactions and adding it in a regularized form. Thus we have

$$\langle \bar{q}_f q_f \rangle_{B,T}^{\text{reg}} = \langle \bar{q}_f q_f \rangle_{B,T} - \langle \bar{q}_f q_f \rangle_{B,T}^{\text{free}} + \langle \bar{q}_f q_f \rangle_{B,T}^{\text{free,reg}} . \quad (14)$$

Notice that “free” condensates are defined keeping the interaction with the magnetic field. In the case of  $\langle \bar{q}_f q_f \rangle_{B,T}^{\text{free,reg}}$  the Matsubara sum can be performed analytically, leading to

$$\langle \bar{q}_f q_f \rangle_{B,T}^{\text{free,reg}} = -\frac{N_c m_c^3}{4\pi^2} \left[ \frac{\ln \Gamma(x_f)}{x_f} - \frac{\ln 2\pi}{2x_f} + 1 - \left( 1 - \frac{1}{2x_f} \right) \ln x_f \right] + \frac{N_c |q_f B|}{\pi} \sum_{k=0}^{\infty} \alpha_k \int \frac{dp}{2\pi} \frac{m_c}{\epsilon_k^f \left[ 1 + \exp(\epsilon_k^f/T) \right]} , \quad (15)$$

where  $x_f = m_c^2/(2|q_f B|)$ ,  $\alpha_k = 2 - \delta_{k0}$  and  $\epsilon_k^f = \sqrt{2k|q_f B| + p^2 + m_c^2}$ . Finally, to make contact with the LQCD results quoted in Ref. [8] we define the quantity

$$\Sigma_{B,T}^f = \frac{2m_c}{S^4} \left[ \langle \bar{q}_f q_f \rangle_{B,T}^{\text{reg}} - \langle \bar{q}_f q_f \rangle_{0,0}^{\text{reg}} \right] + 1 , \quad (16)$$

where the scale  $S$  is given by  $S = (135 \times 86)^{1/2}$  MeV. We also introduce the definitions  $\Delta \Sigma_{B,T}^f = \Sigma_{B,T}^f - \Sigma_{0,T}^f$  and  $\Delta \bar{\Sigma}_{B,T} = (\Delta \Sigma_{B,T}^u + \Delta \Sigma_{B,T}^d)/2$ .

### Numerical results

To obtain the numerical predictions that follow from the above formalism, it is necessary to specify the particular form of the nonlocal form factor. For simplicity, let us consider the often-used Gaussian form  $g(p^2) = \exp(-p^2/\Lambda^2)$ , where the effective scale  $\Lambda$  is an additional parameter of the model. This form factor has the particular advantage that the integral in

Eq. (10) can be performed analytically. One gets

$$M_{\vec{p},k}^{\lambda,f} = m_c + \bar{\sigma} \frac{(1 - |q_f B|/\Lambda^2)^{k+\frac{\lambda_{sf}-1}{2}}}{(1 + |q_f B|/\Lambda^2)^{k+\frac{\lambda_{sf}+1}{2}}} \exp(-\vec{p}^2/\Lambda^2) . \quad (17)$$

Our numerical results for  $T = 0$  are shown in Fig. 1. In the upper panel we quote the model predictions for  $\Delta\bar{\Sigma}_{B,0}$  as a function of  $eB$  for various model parametrizations, while in the lower panel we show the corresponding results for  $\Sigma_{B,0}^u - \Sigma_{B,0}^d$ . LQCD data from Ref. [8] are also displayed in both cases for comparison. Note that the nonlocal model has three parameters, namely,  $m_c$ ,  $G$  and  $\Lambda$ . They have been fixed to reproduce the empirical values of the pion mass and decay constant, and to lead to a certain chosen value of the quark condensate at zero temperature and magnetic field that we identify by  $\Phi_0 \equiv (-\langle\bar{q}_f q_f\rangle_{0,0}^{\text{reg}})^{1/3}$ . Details of this parameter fixing procedure can be found in Ref. [40], where the explicit values of the parameters for  $\Phi_0 = 220$  MeV and 240 MeV are given. As seen in Fig. 1, the predictions for  $\Delta\bar{\Sigma}_{B,0}$  are very similar for all parametrizations considered, and show a very good agreement with LQCD results. In the case of  $\Sigma_{B,0}^u - \Sigma_{B,0}^d$  we see that, although the overall agreement with LQCD calculations is still good, there is a somewhat larger dependence on the model parametrization.

We turn now to our numerical results for the case of finite temperature. In the left panel of Fig. 2 we quote the values obtained for  $\Delta\bar{\Sigma}_{B,T}$  as a function of  $eB$ , for some representative values of the temperature, while in the right panel we show the results for  $(\Sigma_{B,T}^u + \Sigma_{B,T}^d)/2$  as a function of  $T$ , for some selected values of  $eB$ . All these values correspond to the parametrization leading to  $\Phi_0 = 230$  MeV, yet qualitatively similar results are found for the other parametrizations under consideration. The plots in the left panel clearly show that, in contrast to what happens at zero temperature, the quantity  $\Delta\bar{\Sigma}_{B,T}$  does not display a monotonous increase with  $eB$  when one approaches the chiral transition temperature [for this parameter set one has  $T_c(eB = 0) = 129.8$  MeV]. In fact, the curves reach a maximum after which  $\Delta\bar{\Sigma}_{B,T}$  starts to decrease with increasing  $eB$ , implying that the present nonlocal model naturally exhibits the IMC effect found in LQCD. This feature can also be seen from the results displayed in the right panel of Fig. 2. As expected, all curves show a crossover transition from the chiral symmetry broken phase to the (partially) restored one as the temperature increases. However, contrary to what happens e.g. in the standard local NJL model [4–6], it is seen that within the present model the transition temperature decreases as the magnetic field increases. To be more specific, let us define the critical transition

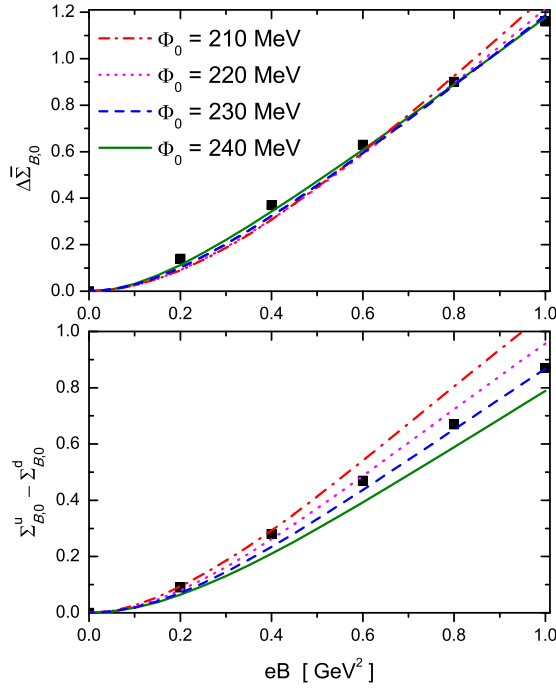


Figure 1: Normalized condensates as functions of the magnetic field at  $T = 0$ . The curves correspond to different model parametrizations identified by  $\Phi_0 = (-\langle \bar{q}_f q_f \rangle_{0,0}^{\text{reg}})^{1/3}$ , for the case of a Gaussian form factor. Full square symbols correspond to LQCD results taken from Ref. [8]. Upper panel: subtracted flavor average; lower panel: flavor difference [see Eq. (16) and the text below].

temperature as the value of  $T$  at which the absolute value of the derivative  $\partial[(\Sigma_{B,T}^u + \Sigma_{B,T}^d)/2]/\partial T$  reaches a maximum. Since, as known from previous analyses [37, 39, 41], the present model is too simple so as to provide realistic values for the critical temperatures even at vanishing external magnetic field, for comparison with LQCD calculations we consider the relative quantity  $T_c(B)/T_c(0)$ . The results corresponding to the previously considered parametrizations are shown in the left panel of Fig. 3, together with LQCD results from Ref. [8]. From the figure it is clearly seen that for magnetic fields beyond  $eB \simeq 0.4 \text{ GeV}^2$  all parameter sets lead to a decrease of the critical temperature when  $eB$  gets increased, i.e. in all cases the IMC effect is observed. In fact, only for the case of  $\Phi_0 = 240 \text{ MeV}$  a slightly opposite behavior is found for lower values of  $eB$ . On the other hand, the strength of the IMC effect is rather sensitive to the parametrization, the best agreement with LQCD being obtained for the parameter set associated with the lowest value of  $\Phi_0$  considered here. Finally, in order to have some estimate of the dependence of our results on the nonlocal



form factor, in the right panel of Fig. 3 we quote the curves corresponding to the relative transition temperatures for the case of a 5-Lorentzian function  $g(p^2) = [1 + (p^2/\Lambda^2)^5]^{-1}$ , often used in the literature. It can be seen that once again the IMC effect is observed for various parametrizations, allowing one to get a reasonably good agreement with LQCD results.

To shed some light on the mechanism that leads to the IMC effect in our model it is worth noticing that the nonlocal form factor turns out to be a function of the external magnetic field. This can be clearly seen from Eq. (10). In addition, it is important to take into account that in nonlocal NJL-like models the form factors play the role of some finite-range gluon-mediated effective interaction. Thus, the magnetic field dependence of the form factor can be understood as originated by the backreaction of the sea quarks on the gluon fields. It is interesting to consider the effective mass for the particular case of a Gaussian form factor, given by Eq. (17). It can be seen that in this case the components of the momentum that are parallel and transverse to the magnetic field become disentangled. While for the 3, 4 components the original exponential form  $\exp(-\vec{p}^2/\Lambda^2)$  is maintained, the 1, 2 (transverse) part leads to a factor given by a ratio of polynomials in  $|q_f B|/\Lambda^2$ , which goes to zero for large  $B$ . One might try to interpret such a factor as a sort of effective magnetic dependent coupling constant, in the line of the analysis carried out e.g. in Ref. [17] in the framework of the local NJL model. However, the analogy is limited by the fact that, contrary to what happens in the case of the local NJL, in our model the so-defined effective coupling is not unique but depends on the Landau level. This important difference prevents a detailed comparison with the particular functional forms used in local NJL analyses. In any case, the qualitative relevant feature is that for any value of  $k$  the strength of the effective coupling decreases as  $eB$  increases. This is analogous to what happens with the  $B$ -dependent coupling constants considered e.g. in Refs. [17, 18], and thus the IMC effect can be understood on these grounds.

### *Summary and outlook*

In this work we have studied the behavior of strongly interacting matter under an external homogeneous magnetic field in the context of nonlocal chiral quark models. These theories are a sort of nonlocal extensions of the local NJL model, intending to represent a step further toward a more realistic modelling of QCD. Considering a Gaussian nonlocal form factor, we have found that at zero temperature the behavior of the quark condensates under

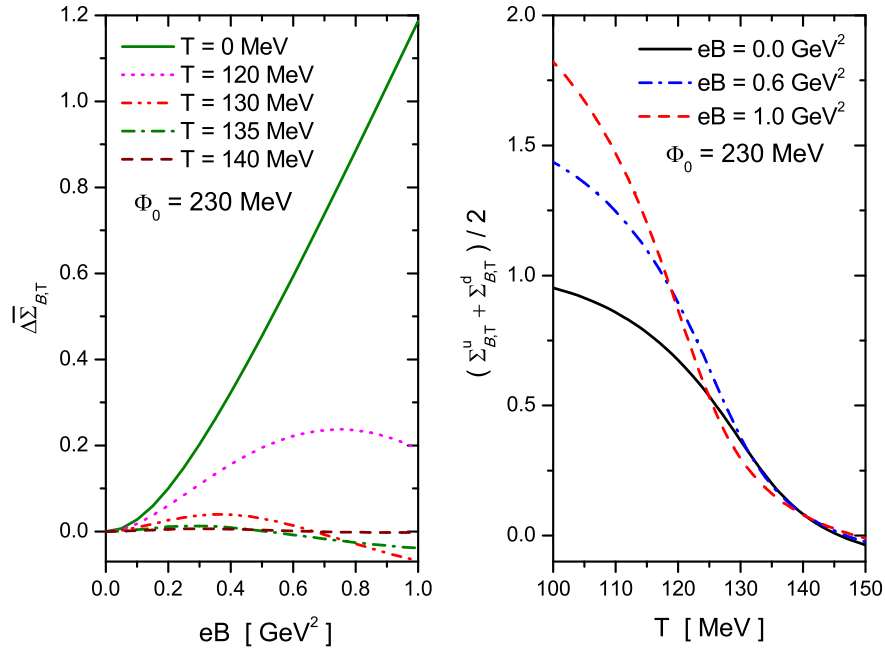


Figure 2: Left: subtracted normalized flavor average condensate as a function of  $eB$  for different representative temperatures. Right: normalized flavor average condensate as a function of the temperature for different representative values of  $eB$ . Results in both panels correspond to  $\Phi_0 = 230$  MeV.

the external field shows the expected magnetic catalysis effect, our predictions being in good quantitative agreement with LQCD results. On the other hand, in contrast to what happens in the standard local NJL model at the mean field level, when the analysis is extended to the case of finite temperature our results show that nonlocal models naturally lead to the Inverse Magnetic Catalysis effect. It is worth stressing that in these models the current-current couplings turn out to be dependent on the temperature and the magnetic field through the nonlocal form factors, which in principle follow from some finite-range gluon-mediated effective interaction. Our results indicate that this scheme seems to capture the main features of more sophisticated approaches to the QCD dynamics in the presence of external magnetic fields, in which IMC is observed. We have also analyzed the numerical results obtained for other form factor shapes often considered in the literature (see e.g. Ref. [40]). For comparison we have quoted in this work the results for the relative critical temperatures corresponding to a 5-Lorentzian form factor, which also show the presence

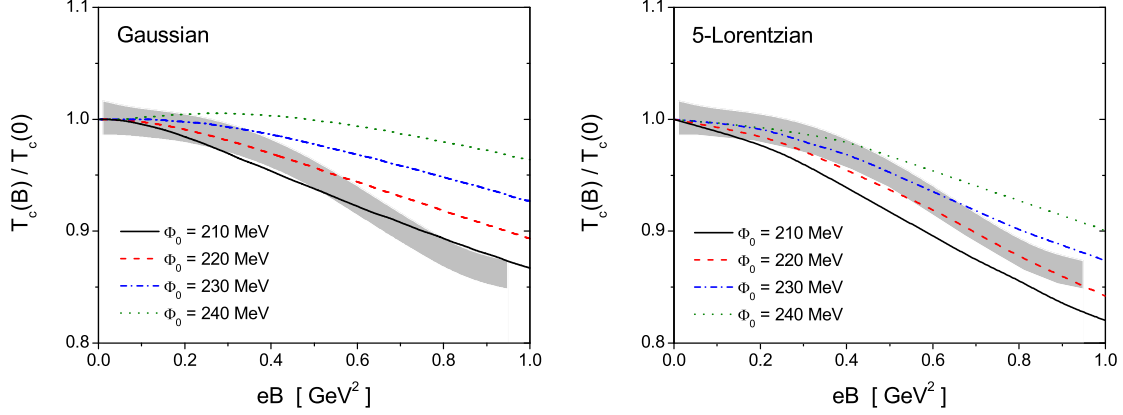


Figure 3: Normalized chiral restoration temperatures as functions of  $eB$  for various model parametrizations. For comparison, LQCD results of Ref. [8] are indicated by the gray band. Left and right panel correspond to Gaussian and 5-Lorentzian form factors, respectively.

of the IMC effect. A further analysis of the predictions arising from other possible form factors, together with a more extended presentation of the formalism, will be provided in a forthcoming article [49]. It is also worth noticing that, as a first step in this research line, we have considered here a simple version of nonlocal models in which e.g. we have not incorporated interactions leading to quark wave function renormalization nor the coupling of fermions to the Polyakov loop. It is clear that the inclusion of these interactions is important to provide a more realistic description of strong interaction thermodynamics [41, 42]. We expect to report progresses in this direction in the near future.

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- [1] R. C. Duncan and C. Thompson, *Astrophys. J.* **392**, L9 (1992); C. Kouveliotou et al., *Nature* **393**, 235 (1998).
- [2] D. E. Kharzeev, L. D. McLerran and H. J. Warringa, *Nucl. Phys. A* **803**, 227 (2008); V. Skokov, A. Y. Illarionov, and V. Toneev, *Int. J. Mod. Phys. A* **24**, 5925 (2009); V. Voronyuk, V. Toneev, W. Cassing, E. Bratkovskaya, V. Konchakovski, and S. Voloshin, *Phys. Rev. C* **83**, 054911 (2011).
- [3] T. Vachaspati, *Phys. Lett.* **B265**, 258 (1991); K. Enqvist and P. Olesen, *Phys. Lett.* **B319**, 178 (1993).
- [4] D. E. Kharzeev, K. Landsteiner, A. Schmitt and H. U. Yee, *Lect. Notes Phys.* **871**, 1 (2013).
- [5] J. O. Andersen, W. R. Naylor and A. Tranberg, *Rev. Mod. Phys.* **88**, 025001 (2016).
- [6] V. A. Miransky and I. A. Shovkovy, *Phys. Rept.* **576**, 1 (2015).
- [7] G. S. Bali, F. Bruckmann, G. Endrodi, Z. Fodor, S. D. Katz, S. Krieg, A. Schafer and K. K. Szabo, *JHEP* **1202**, 044 (2012).
- [8] G. S. Bali, F. Bruckmann, G. Endrodi, Z. Fodor, S. D. Katz and A. Schafer, *Phys. Rev. D* **86**, 071502 (2012).
- [9] V. Skokov, *Phys. Rev. D* **85**, 034026 (2012).
- [10] E. S. Fraga, J. Noronha and L. F. Palhares, *Phys. Rev. D* **87**, 114014 (2013).
- [11] F. Bruckmann, G. Endrodi and T. G. Kovacs, *JHEP* **1304**, 112 (2013).
- [12] G. S. Bali, F. Bruckmann, G. Endrodi, F. Gruber and A. Schaefer, *JHEP* **1304**, 130 (2013).
- [13] K. Fukushima and Y. Hidaka, *Phys. Rev. Lett.* **110**, 031601 (2013).
- [14] J. Chao, P. Chu and M. Huang, *Phys. Rev. D* **88**, 054009 (2013).
- [15] E. S. Fraga, B. W. Mintz and J. Schaffner-Bielich, *Phys. Lett.* **B731**, 154 (2014).
- [16] M. Ferreira, P. Costa, D. P. Menezes, C. Providência and N.N. Scoccola, *Phys. Rev. D* **89**, 016002 (2014);
- [17] M. Ferreira, P. Costa, O. Lourenço, T. Frederico and C. Providência, *Phys. Rev. D* **89**, 116011 (2014).
- [18] R. L. S. Farias, K. P. Gomes, G. I. Krein and M. B. Pinto, *Phys. Rev. C* **90**, 025203 (2014).
- [19] A. Ayala, M. Loewe, A. J. Mizher and R. Zamora, *Phys. Rev. D* **90**, 036001 (2014).

- [20] A. Ayala, M. Loewe and R. Zamora, Phys. Rev. D **91**, 016002 (2015); A. Ayala, C. A. Dominguez, L. A. Hernandez, M. Loewe and R. Zamora, Phys. Rev. D **92**, 096011 (2015); Addendum: [Phys. Rev. D **92**, 119905 (2015)].
- [21] S. Fayazbakhsh and N. Sadooghi, Phys. Rev. D **90**, 105030 (2014).
- [22] J. O. Andersen, W. R. Naylor and A. Tranberg, JHEP **1502**, 042 (2015).
- [23] N. Mueller and J. M. Pawłowski, Phys. Rev. D **91**, 116010 (2015).
- [24] A. Ayala, J. J. Cobos-Martínez, M. Loewe, M. E. Tejeda-Yeomans and R. Zamora, Phys. Rev. D **91**, 016007 (2015).
- [25] E. J. Ferrer, V. de la Incera and X. J. Wen, Phys. Rev. D **91**, 054006 (2015).
- [26] J. Braun, W. A. Mian and S. Rechenberger, Phys. Lett. **B755**, 265 (2016).
- [27] R. Rougemont, R. Critelli and J. Noronha, Phys. Rev. D **93**, 045013 (2016).
- [28] A. Ayala, C. A. Dominguez, L. A. Hernandez, M. Loewe and R. Zamora, Phys. Lett. **B759**, 99 (2016).
- [29] S. Mao, Phys. Lett. **B758**, 195 (2016).
- [30] T. Schafer and E. V. Shuryak, Rev. Mod. Phys. **70**, 323 (1998).
- [31] C. D. Roberts and A. G. Williams, Prog. Part. Nucl. Phys. **33**, 477 (1994); C. D. Roberts and S. M. Schmidt, Prog. Part. Nucl. Phys. **45**, S1 (2000).
- [32] S. Noguera and N. N. Scoccola, Phys. Rev. D **78**, 114002 (2008).
- [33] G. Ripka, *Quarks Bound by Chiral Fields*, (Oxford University, New York, 1997).
- [34] R.D. Bowler and M.C. Birse, Nucl. Phys. A **582**, 655 (1995); R.S. Plant and M.C. Birse, Nucl. Phys. A **628**, 607 (1998).
- [35] S. M. Schmidt, D. Blaschke and Y. L. Kalinovsky, Phys. Rev. C **50**, 435 (1994).
- [36] B. Golli, W. Broniowski and G. Ripka, Phys. Lett. B **437**, 24 (1998); W. Broniowski, B. Golli and G. Ripka, Nucl. Phys. A **703**, 667 (2002).
- [37] I. General, D. Gomez Dumm and N. N. Scoccola, Phys. Lett. **B506**, 267 (2001); D. Gomez Dumm and N.N. Scoccola, Phys. Rev. D **65**, 074021 (2002).
- [38] A. Scarpettini, D. Gomez Dumm and N.N. Scoccola, Phys. Rev. D **69**, 114018 (2004).
- [39] D. Gomez Dumm and N. N. Scoccola, Phys. Rev. C **72**, 014909 (2005).
- [40] D. Gomez Dumm, A. G. Grunfeld and N.N. Scoccola, Phys. Rev. D **74**, 054026 (2006).
- [41] G. A. Contrera, D. Gomez Dumm and N. N. Scoccola, Phys. Lett. **B661**, 113 (2008); G. A. Contrera, M. Orsaria and N. N. Scoccola, Phys. Rev. D **82**, 054026 (2010); J. P. Car-

- lomagno, D. Gómez Dumm and N. N. Scoccola, Phys. Rev. D **88**, 074034 (2013).
- [42] T. Hell, S. Roessner, M. Cristoforetti and W. Weise, Phys. Rev. D **79**, 014022 (2009); T. Hell, S. Rossner, M. Cristoforetti and W. Weise, Phys. Rev. D **81**, 074034 (2010).
- [43] D. Gomez Dumm, S. Noguera and N.N. Scoccola, Phys. Lett. B **698**, 236 (2011); Phys. Rev. D **86**, 074020 (2012).
- [44] K. Kashiwa, Phys. Rev. D **83**, 117901 (2011).
- [45] V. I. Ritus, Sov. Phys. JETP **48**, 788 (1978).
- [46] C. N. Leung, Y. J. Ng and A. W. Ackley, Phys. Rev. D **54**, 4181 (1996).
- [47] P. Watson and H. Reinhardt, Phys. Rev. D **89**, 045008 (2014).
- [48] N. Mueller, J. A. Bonnet and C. S. Fischer, Phys. Rev. D **89**, 094023 (2014).
- [49] V.P. Pagura, D. Gomez Dumm, S. Noguera and N.N. Scoccola, in preparation.